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R-symmetric axion/natural inflation in supergravity via deformed moduli dynamics

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ABSTRACT

We construct a natural inflation model in supergravity where the inflaton is identified with a modulus field possessing a shift symmetry. The superpotential for the inflaton is generated by meson condensation due to strong dynamics with deformed moduli constraints. In contrast to models based on gaugino condensation, the inflaton potential is generated without R -symmetry breaking and hence does not depend on the gravitino mass. Thus, our model is compatible with low scale supersymmetry.

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1. Introduction

Slow-roll inflation [1,2] is now a standard paradigm in the modern cosmology. It not only solves the flatness problem and the horizon problem [3,4], but it also explains the origin of the large scale structure of the universe [5–9]. This paradigm has been supported by precise measurements of the cosmic microwave background (CMB) [10–12].

After the announcement by the BICEP2 experiment on the B-mode polarization [13], models with larger inflaton field values than the Planck scale are drawing much attention due to the so-called Lyth bound [14].¹ Such a large field value seems inconsistent with the conventional view of the field theoretic description as an effective theory which is believed to be at the best given by a series expansion of fields with higher dimensional operators suppressed by the Planck scale. In other words, in large field inflation models, any higher dimensional terms of the inflaton potential should be somehow under control.

The best way to understand such strict control on the inflaton potential would be a shift symmetry of the inflaton [16]. Interestingly, such a candidate of the inflaton with a shift symmetry is often provided in string theories as a modulus [17]. We refer to the modulus as an axion, although it is not the QCD axion which solves the strong CP problem [18–20]. Once we identify the axion

as the inflaton, the next task is to generate a potential of the axion. As a caveat, in the situation where the shift symmetry holds at the tree level and is broken by quantum effects, as is often the case with axions in superstring theories, the superpotential of the axion, and hence, the axion potential, is generated only by non-perturbative effects [21]. Thus, model construction often requires strong gauge dynamics to generate the axion potential.

Along this line, natural inflation models in supergravity have been constructed [27–33].² As a common feature of these models, the axion potential originates from gaugino condensation in strongly coupled gauge theories. As a result, the energy scale of the axion potential is proportional to the scale of R symmetry breaking, i.e., the gravitino mass. Thus, to explain the magnitude of cosmic perturbations, the gravitino mass is required to be as large as 10^{13} GeV, which is incompatible with low scale supersymmetry breaking.³

In this letter, we propose to make use of meson condensation by strong dynamics with deformed moduli constraints to generate the superpotential of the axion/inflaton field.⁴ As we will show, the model possesses an R -symmetry and the inflaton potential does not depend on the gravitino mass. Thus, our model is compatible with low energy supersymmetry breaking.

² For inflation models other than natural inflation where inflaton potentials are generated dynamically, see Refs. [22–26].

³ In Refs. [34,35], natural inflation models consistent with low scale supersymmetry breaking are proposed, although the shift symmetry breaking is simply given by tree-level superpotentials.

⁴ The idea of generating axion potential by the meson condensation is suggested in Ref. [36].

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¹ Models with large inflaton field value are free from the initial condition problem [15].

2. Inflaton potential from meson condensation

2.1. Dynamical sector

Let us begin with a brief review on a supersymmetric $SP(N_c)$ gauge theory with $2(N_c + 1)$ chiral superfields in the fundamental representation, Q^i ($i = 1 \cdots 2(N_c + 1)$).⁵ The vacuum structure of classical flat directions, i.e. $(N_c + 1)(2N_c + 1)$ meson fields,

$$M^{ij} \propto Q^i Q^j, \quad (1)$$

is deformed non-perturbatively. The vacuum expectation values (VEVs) of the meson fields obey the so-called deformed quantum moduli constraint [37];

$$Pf^{(N_c+1)}(M^{ij}) = \Lambda^{N_c+1}. \quad (2)$$

Here, Λ denotes the dynamical scale of the $SP(N_c)$ gauge interaction and $Pf(\cdots)$ denotes the Pfaffian.⁶ We have normalized the meson fields M^{ij} so that they have a mass dimension one. As is clear from Eq. (2), some of the mesons condensate at the vacuum.

2.2. Axion

Next, let us introduce an axion chiral multiplet T which couples to the above gauge dynamics via the gauge kinetic function. Later on, we will identify the imaginary part of the axion multiplet T with the inflaton. To be concrete, we assume that the Kähler potential of the axion multiplet is given by⁷

$$K = K(T + T^\dagger) = \frac{1}{2}(T + T^\dagger)^2 + \cdots, \quad (3)$$

where the ellipses denotes higher dimensional terms. Here, we have assumed that the Kähler potential has a shift symmetry, $T \rightarrow T + i\alpha$, with α being a real number. We also assume that the axion multiplet appears in the gauge kinetic function of the $SP(N_c)$ gauge multiplet,

$$\mathcal{L}_{\text{gauge}} = \frac{1}{4} \int d^2\theta \left(\frac{1}{g^2} + \frac{T}{8\pi^2 f_a} \right) W^\alpha W_\alpha + \text{h.c.}, \quad (4)$$

where a dimensionful constant f_a denotes the “decay constant” which depends on the origin of the axion multiplet. We assume that this coupling is the dominant contribution to the shift symmetry breaking of the axion.

In our argument, instead of specifying the origin of the axion multiplet, we simply assume that the value of f_a is at around the so-called string scale, i.e. $M_{\text{str}} \simeq 10^{17}$ GeV, which is expected in the case of string axions [17]. Through the coupling to the gauge kinetic term, the shift symmetry is broken by the non-perturbative effects of the $SP(N_c)$ dynamics.

2.3. STEP1

In the presence of the axion multiplet in the kinetic function, the effective dynamical scale depends on the axion field, i.e.,

$$\Lambda_{\text{eff}}(T) = \Lambda \exp \left[-\frac{1}{2(N_c + 1)} \frac{T}{f_a} \right]. \quad (5)$$

Accordingly, the above meson condensation in Eq. (2) also depends on the axion multiplet, i.e.

$$Pf^{(N_c+1)}(M^{ij}) = \Lambda_{\text{eff}}^{N_c+1}(T). \quad (6)$$

It should be emphasized here that mere condensation of the mesons does not lead to a non-trivial potential of the axion multiplet, although the meson condensation scale depends on the axion multiplet. This feature should be contrasted with the axion potential generation via the gaugino condensation, where the condensation leads to a non-trivial potential of the axion multiplet.

2.4. STEP2

To generate a non-trivial axion potential, let us introduce $(N_c + 1)(2N_c + 1)$ singlet fields, $X^{ij} = -X^{ji}$, which couple to the fundamental fields Q^i in the same way with the model of dynamical supersymmetry breaking developed in Refs. [38,39];

$$W = \sum_{i>j, k>l} \lambda_{ij,kl} X^{ij} Q^k Q^l. \quad (7)$$

To make our analysis simple, we hereafter assume that the above superpotential possesses a global $SP(2(N_c + 1))$ symmetry out of the maximal flavor $SU(2(N_c + 1))$ symmetry, and that the $SP(2(N_c + 1))$ singlet direction, $X^{ij} \propto J^{ij}$, has the smallest coupling to the quarks, i.e.,⁸

$$\lambda_{ij,kl} = \lambda'' J_{ik} J_{jl} + \left(\lambda' - \frac{\lambda''}{2(N_c + 1)} \right) J_{ij} J_{kl} \quad (|\lambda'| < |\lambda''|). \quad (8)$$

Below the dynamical scale, the tree-level interactions lead to effective couplings between the mesons and the singlets,

$$W_{\text{eff}} \simeq \sum_{i>j, k>l} \lambda_{ij,kl} \Lambda_{\text{eff}}(T) X^{ij} M^{kl}, \quad (9)$$

where the mesons are subject to the deformed constraint in Eq. (6).⁹ In this effective theory, we see that all the meson fields and the singlets get massive at around the VEVs of the mesons,

$$M^{ij} = \Lambda_{\text{eff}}(T) \times J^{ij}, \quad (10)$$

except for the singlet which corresponds to the global $SP(2(N_c + 1))$ singlet.¹⁰ By inserting this solution to the effective potential, we obtain the effective superpotential of the remaining singlet field,

$$W_{\text{eff}} \simeq \lambda \Lambda_{\text{eff}}(T)^2 X \simeq \lambda \Lambda^2 e^{-\frac{1}{2(N_c+1)} \frac{T}{f_a}} X, \quad (11)$$

after integrating out other heavier mesons and singlets.¹¹ Here, we have defined

$$X = \frac{1}{\sqrt{N_c + 1}} \sum_{i>j} J^{ij} X_{ij}, \quad (12)$$

$$\lambda = 2\lambda' (N_c + 1)^{3/2}. \quad (13)$$

As a result, we find that the supersymmetry is broken for a given value of T , which leads to a non-trivial potential of the axion field,

⁸ The following arguments can be extended to generic cases as done in Ref. [26].

⁹ We may consider the deformed moduli constraint as a consequence of equations of motions of heavy states such as glueball supermultiplet of $SP(N_c)$. Following arguments are not significantly altered even when we treat the deformed moduli constraints as the equation of motion of heavy states.

¹⁰ One of the meson obtains a mass of $O(\Lambda_{\text{eff}})$ due to the deformed constraint.

¹¹ Mixing between the axion and the mesons is suppressed by $\Lambda_{\text{eff}}/((N_c + 1)f_a)$ and hence negligible [36].

⁵ In our convention, $SP(1)$ is equivalent to $SU(2)$.

⁶ We define the Pfaffian of a $2n \times 2n$ antisymmetric matrix, $Pf^{(n)}$, so that the symplectic form J , where $J = \mathbb{1}_n \otimes i\sigma_2$ with $\mathbb{1}_n$ being the $n \times n$ unit matrix and σ_2 being the second Pauli matrix, satisfies $Pf^{(n)}(J) = 1$.

⁷ Here, we have chosen the origin of T so that the Kähler potential does not have a linear term $T + T^\dagger$.

$$V_{\text{eff}} \simeq |\lambda|^2 \Lambda^4 e^{-\frac{1}{(N_c+1)} \frac{T+T^\dagger}{f_a}}, \quad (14)$$

where we have set $X = 0$.¹² Unfortunately, however, the imaginary part of the scalar component of T , the axion field, remains flat, and hence, this dynamics does not lead to the model of natural inflation.

2.5. STEP3

The above failure can be traced back to the remaining shift-symmetry in the effective potential in Eq. (11) under which X rotates to absorb the shift of T .¹³ Therefore, to generate non-trivial potential for the imaginary part of the axion, we are lead to add a linear term of X which breaks the remaining shift-symmetry explicitly,

$$\Delta W = -\mu^2 X, \quad (15)$$

where μ is a dimensionful parameter.¹⁴ We note that this term is consistent with the R -symmetry which we discuss later. In the presence of the breaking term, the above dynamics leads to the effective potential,

$$W = \lambda \Lambda^2 \left(e^{\frac{1}{(N_c+1)} \frac{T}{f_a}} - \tilde{\mu}^2 \right) X, \quad (16)$$

$$\tilde{\mu}^2 = \frac{\mu^2}{\lambda \Lambda^2}. \quad (17)$$

In the followings, we take a phase convention of X so that $\tilde{\mu}$ is real and positive valued. As a result, we obtain an axion potential,

$$V_{\text{eff}} \simeq |\lambda|^2 \Lambda^4 \left(e^{\sqrt{2} \frac{\tau}{(N_c+1) f_a}} + \tilde{\mu}^4 \right) - 2\tilde{\mu}^2 e^{\frac{\tau}{\sqrt{2}(N_c+1) f_a}} \cos \left[\frac{\phi}{\sqrt{2}(N_c+1) f_a} \right], \quad (18)$$

which lifts up the imaginary part of the axion field. In the above expression, we have decomposed the axion field into

$$T = \frac{1}{\sqrt{2}}(\tau + i\phi). \quad (19)$$

It should be noted that unlike the model of dynamical supersymmetry breaking model in [38,39], the model does not break supersymmetry spontaneously due to the presence of T , where the supersymmetry vacuum is at

$$e^{\frac{\tau}{\sqrt{2}(N_c+1) f_a}} \simeq \tilde{\mu}^2, \quad \phi = 0. \quad (20)$$

At around this vacuum, both the axion and its real field counterpart obtain the same mass,

¹² Here, it should be noted that the scalar component of X is stabilized to $X = 0$ due to a large positive mass of the scalar component, $\Delta m_X^2 \sim \lambda^4 H_{\text{inf}}^2 M_{\text{Pl}}^2 / \Lambda_{\text{eff}}^2$, generated by perturbative corrections [40]. In the limit of a small dynamical scale, the mass of the scalar component is far larger than the Hubble scale and the scalar component decouples during inflation. Thus, X can be identified with a nilpotent chiral superfield discussed in Ref. [41]. In our model, however, Δm_X^2 vanishes and the scalar component becomes light after inflation. The chiral multiplet X becomes a mass partner of the inflaton and is relevant for the decay of the inflaton.

¹³ Accordingly, the fundamental fields also rotate under the remaining symmetry which makes the original shift symmetry free of the anomaly, and hence, one linear combination of the phases remains as a massless axion.

¹⁴ If we extend the definition of the shift symmetry so that X rotates non-trivially, we may add a term $\Delta W' = e^{cT} X$ with an appropriate coefficient c , which is in general broken by the anomaly of the $SP(N_c)$ gauge interaction. With such a term, we obtain a different inflaton potential, although we do not pursue this possibility in this paper.

$$m^2 = \frac{\mu^4}{(N_c+1)^2 f_a^2}. \quad (21)$$

It should be also noted that the resultant scalar potential does not show the runaway behavior as seen in Eq. (14).

Now, let us assume that the real part of the axion field is fixed to its supersymmetric vacuum value, while allowing the axion field being away from its vacuum value, i.e. $\phi \neq 0$. In this case, the superpotential Eq. (17) is reduced to

$$W = \sqrt{2} \mu^2 \left(e^{\frac{1}{(N_c+1)} \frac{i\phi}{\sqrt{2} f_a}} - 1 \right) X, \quad (22)$$

where ϕ should be understood as not a chiral field but a constant. It should be noted that our model has the same structure as the “Model 1” in Ref. [34], and hence, our model provides an ultraviolet completion to their model.

As a result, along the lines of the chaotic inflation model with shift symmetry in [16,34], the axion field obtains a non-trivial potential through the F -term contribution of X which leads to

$$V_{\text{eff}} \simeq 2\mu^4 \left(1 - \cos \left[\frac{\phi}{f_{\text{eff}}} \right] \right). \quad (23)$$

Here, we have defined an effective decay constant,

$$f_{\text{eff}} = \sqrt{2}(N_c+1) f_a. \quad (24)$$

It should be noted that the effective decay constant is required to be larger than the Planck scale to satisfy the slow-roll conditions in natural inflation. For $f_a = O(M_{\text{str}}) = O(10^{17})$ GeV, the effective decay constant is larger than the Planck scale if $N_c = O(10)$.¹⁵

In this way, we find that the model with meson condensation leads to the inflaton potential which is appropriate for natural inflation. For recent discussion on the consistency of natural inflation with CMB data, we refer e.g. Ref. [43].

2.6. Required tuning

We clarify how feasible it is to assume that the real part of τ is fixed to a desirable position in Eq. (20). For that purpose, let us first estimate the mass of the real part of the axion around the field value in Eq. (20), which is given by,

$$m_\tau^2 = \frac{4\mu^4}{f_{\text{eff}}^2} \ll \frac{4\mu^4}{f_{\text{eff}}^2} \frac{\phi^2}{M_{\text{Pl}}^2} \simeq H_{\text{inf}}^2, \quad (25)$$

where H_{inf} denotes the Hubble parameter during inflation. Thus, the real part of the axion is not fixed by the superpotential coupling to X , and hence, we need to have the axion coupling to X in the Kähler potential which is in general given by,

$$\Delta K = \frac{X^\dagger X}{M_{\text{Pl}}^2} (\sqrt{2} c_1 M_{\text{PL}} (T + T^\dagger) + c_2 (T + T^\dagger)^2 / 2 + \dots), \quad (26)$$

where $c_{1,2}$ are $O(1)$ coefficients. With these terms, the real part of the axion field is fixed to

$$\tau_* \simeq \frac{c_1}{1 - c_2} M_{\text{PL}}, \quad (27)$$

where we have assumed $|c_1| \ll 1$, for simplicity. In general, this field value is expected to be far away from the vacuum position in Eq. (20).

¹⁵ Enhancement of the effective decay constant in the inflaton potential by a large N_c is pointed out in Refs. [32,33,42]. In the view point of the “Phase Locking Mechanism” proposed in Refs. [35], the enhanced decay constant is understood by hierarchical charges between the phase of X and ϕ under the remaining shift symmetry discussed at the beginning of this subsection, and the breaking of the remaining shift symmetry by the superpotential term in Eq. (15).

If the real part of the axion field is fixed at far away from the vacuum position, the axion stays at τ_* , and never goes back to the vacuum position after inflation since the effective mass of the real part of the axion around τ_* is much larger than the one in Eq. (25). Hence, inflation never ends due to the non-vanishing potential energy at τ_* even for $\phi = 0$, i.e.

$$V_{\text{eff}} \simeq |\lambda|^2 \Lambda^4 \left(e^{\frac{\tau}{\sqrt{2(N_c+1)}f_a}} - \tilde{\mu}^2 \right)^2 \neq 0. \quad (28)$$

Thus, in order to avoid this problem, we need to tune the value of μ , so that

$$e^{\frac{\tau_*}{\sqrt{2(N_c+1)}f_a}} = \tilde{\mu}^2 (1 + \delta) \quad (\delta \ll 1). \quad (29)$$

With this tuning, the inflaton potential along τ_* is given by

$$V_{\text{eff}} \simeq 2\mu^4 (1 + \delta) \left(1 + \frac{\delta^2}{2} - \cos \left[\frac{\phi}{\sqrt{2(N_c+1)}f_a} \right] \right). \quad (30)$$

By remembering that the mass of τ around τ_* is given by,

$$m_{\tau_*}^2(\phi) \simeq \frac{(1 - c_2)V_{\text{eff}}(\phi)}{M_{\text{PL}}^2}, \quad (31)$$

we find that the axion field goes back to the vacuum position well after inflation, i.e. $\phi \simeq 0$ as long as

$$m_{\tau_*}^2 > m_{\tau_*}^2(\phi \simeq 0) \simeq \frac{(1 - c_2)\mu^4 \delta^2}{M_{\text{PL}}^2}. \quad (32)$$

To satisfy the above condition, we find that we need tuning between parameters,

$$\delta < \frac{4M_{\text{PL}}^2}{(1 - c_2)f_{\text{eff}}^2}. \quad (33)$$

2.7. Energy scale of parameters

Here, we summarize the energy scale of parameters in our model. The observed magnitude and the spectral index of the curvature perturbation require that [43]

$$f_{\text{eff}} > 5M_{\text{Pl}}, \quad \frac{\mu^2}{f_{\text{eff}}} \sim 10^{-5}M_{\text{Pl}}. \quad (34)$$

The dynamical scale of the $SP(N_c + 1)$ gauge theory during inflation is

$$\Lambda(\tau_*)^2 \simeq \lambda^{-1} \mu^2 \sim 10^{-5} \lambda^{-1} f_{\text{eff}} M_{\text{Pl}}. \quad (35)$$

The Hubble scale during inflation is given by

$$H_{\text{inf}}^2 \sim \frac{\mu^4 \phi^2}{M_{\text{Pl}}^2 f_{\text{eff}}^2} \sim 10^{-10} \phi^2. \quad (36)$$

Note that the dynamical scale is larger than the Hubble scale as long as $\phi^2 < 10^5 \lambda^{-1} f_{\text{eff}} M_{\text{Pl}}$. Thus, the Hubble induced masses of quarks Q^i are negligible and hence does not affect the gauge dynamics.

2.8. R -symmetry

Finally, we note that the R -symmetry is preserved in our model. The R -charge assignment is $X(2)$, $Q^i(0)$ and $T(0)$. The R -symmetry is free from the gauge anomaly of the $SP(N_c)$, and hence not explicitly broken by the strong dynamics of the $SP(N_c)$ gauge theory. Also, since the scalar component of X is fixed to its origin, the R -symmetry is also not spontaneously broken. Thus, the

inflaton sector does not break the R -symmetry, and hence the inflation scale is not related with the gravitino mass. Our model is compatible with low scale supersymmetry.

We stress that the R -symmetry is important for stable inflation dynamics. If the R -symmetry is broken during inflation, the negative contribution to the inflaton potential is significant and the inflaton may be destabilized toward far from the origin. In our model, since the R -symmetry is preserved during inflation, the negative contribution is absent.

We have made use of meson condensation to generate the inflaton potential. As is pointed out in Ref. [44], the mechanism can be applied to moduli fixing. Since moduli are fixed in an R invariant way, masses of moduli can be far larger than the gravitino mass. Thus, moduli fixing by meson condensation is free from destabilization of moduli during inflation by Hubble induced potentials.

3. Summary and discussion

In this letter, we have proposed a natural inflation model in supergravity where the axion potential is generated by meson condensation due to strong dynamics with deformed moduli constraints. In contrast to models based on gaugino condensation, our model possesses an unbroken R -symmetry and hence the inflaton potential does not depend on the gravitino mass. Thus, our model is compatible with low scale supersymmetry.

In the above analysis, we have assumed one axion field. It is easy to extend our model to multi-axion cases. For example, let us consider two axions T and S . We couple them to two gauge theories via gauge kinetic functions and assume that gauge theories are in meson condensation phases. By fixing mesons in the same way as the above analysis, we obtain the effective super potential,

$$W_{\text{eff}} = X \Lambda^2 \left(\exp \left[\frac{T}{f_T} + \frac{S}{f_S} \right] - \tilde{\mu}^2 \right) + X' \Lambda'^2 \left(\exp \left[\frac{T}{f'_T} + \frac{S}{f'_S} \right] - \tilde{\mu}'^2 \right), \quad (37)$$

where X and X' are singlet fields corresponding to that in Eq. (11) for two gauge theories, and $\Lambda^{(\prime)}$, $f_T^{(\prime)}$, $f_S^{(\prime)}$ and $\tilde{\mu}^{(\prime)}$ are constants. If $f_T/f_S \simeq f'_T/f'_S$, a linear combination of T and S works as an inflaton with an effective decay constant much larger than $f_T^{(\prime)}$ and $f_S^{(\prime)}$ [29].

We have assumed the global $SP(2(N_c + 1))$ symmetry to simplify our analysis. Without the symmetry, the VEVs of mesons are not given by Eq. (10), but generic ones which depend on constants λ s and μ^2 s. After integrating out heavy mesons and singlets, the effective superpotential is given by Eq. (11), but λ , $\tilde{\mu}^2$ and Λ in general depends on T . As a result, the inflaton potential is not given by a simple cosine form. It is interesting if deviation from the cosine form is observed.

Let us comment on decay of the inflaton. The inflaton does not possess any charges under some linearly realized symmetry. Thus, the inflaton in general decays into standard model particles through its linear terms in the Kähler potential or gauge kinetic functions [45,46].

The inflaton also decays into supersymmetry breaking sector fields, which may lead to the overproduction of gravitinos [45–50]. The overproduction can be avoided if masses of supersymmetry breaking sector fields are large enough, so that the decay mode is kinematically forbidden [51].¹⁶ As is discussed in Ref. [52], the supersymmetry breaking scale may have a lower bound. It may be

¹⁶ The inflaton also decays into a pair of gravitinos through the mixing between the inflaton and a scalar component of the supersymmetry breaking field. This de-

interesting that the large supersymmetry breaking scale assumed in the pure gravity mediation [53–55] is naturally explained in this way.

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